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Towards a viable grand unified model with $M_G \sim M_{\text{string}}$ and $M_I \sim 10^{12}$ GeV

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Abstract

We present a model based on the gauge group $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{SU}(4)_C$ with gauge couplings that are found to be unified at a scale M_G near the string unification scale. The model breaks to the minimal supersymmetric standard model at a scale $M_I \sim 10^{12}$ GeV, which is instrumental in producing a neutrino in a mass range that can serve as hot dark matter and this scale can also solve the strong CP problem via the Peccei-Quinn (PQ) mechanism with an invisible harmless axion. We show how this model can accommodate low and high values of $\tan\beta$ and "exotic" representations that often occur in string derived models. We show that this model has lepton flavor violation which can lead to processes which are one or two orders of magnitude below the current experimental limits.

The conventional scale of supersymmetric grand unification is taken to be $M_G \sim 2 \cdot 10^{16}$ GeV, because this is where the MSSM gauge couplings are found to converge if one assumes a "dessert" between about 1 TeV and and that scale. However, in superstring theory the unification point is not a free parameter but is predicted to be a function of the gauge coupling at that scale in the $\overline{\mathrm{MS}}$ scheme as follows [1]:

$$M_{\mathrm{string}} \approx 7 g_{\mathrm{string}} \cdot 10^{17}\,\mathrm{GeV}\,, \tag{1}$$

which predicts M_{string} to be approximately 25 times greater than the conventional value of M_G. Stringy threshold effects have not yet proven to be at all useful in closing the gap between M_G and M_{string} in any realistic string model [2], and neither have weak to TeV scale MSSM thresholds. At present, it is not clear if string grand unified theories (GUTs) or models that have non-standard Kac-moody levels and hence non-conventional hypercharge normalizations may one day be able to rectify the situation [2]. However, it has been shown that extra non-MSSM matter that appears in some realistic string models can lead to a successful raising of the unification scale [2]. One obvious and attractive approach to adding extra matter, "populating the dessert," would be to add an intermediate gauge symmetry. Such realistic string models have been built for cases where the intermediate symmetry is the flipped SU(5)× U(1) [3] or SO(4)× SO(6)∼ $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{SU}(4)_C$ [4–7], and sometimes the field content of these models have been found to alleviate the discrepancy between the string and gauge unification scales [6–8]. In this letter, we shall investigate what field content and additional constraints may be required to have a model with gauge unification at the string scale, that has $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{SU}(4)_C$ breaking to the MSSM at an intermediate scale of $\sim 10^{12}$ GeV, that allows a PQ symmetry also broken $\sim 10^{12}$ GeV with a weakly coupled axion, and where the scale of gauge symmetry breaking and the PQ symmetry breaking are determined by a single parameter in the superpotential. We will also discuss the constraint on the field content necessary to obtain, if one desires, a low value for $\tan\beta$.

An intermediate $\mathrm{SU}(2)_R \times \mathrm{SU}(4)_C$ breaking scale of order 10^{12} GeV is very attractive for two reasons: (1) if the B-L gauge symmetry is broken at around 10^{11}-10^{12} GeV, one

can easily get a neutrino mass in the interesting range of of about 3-10 eV, making it a candidate for the hot dark matter [9], and (2) if the strong CP problem is solved via the Peccei-Quinn (PQ) mechanism, this PQ symmetry is required to be broken approximately within the above window so that the axion has properties which are consistent with the lack of observation up to now and the cosmological constraints [10]. As for the scale at which the hypothetical PQ-symmetry is broken, perhaps the most elegant possibility is if it is tied in with the breaking of an intermediate gauge symmetry, so that there is only one scale between the weak and string scale to be explained. To obtain the τ-neutrino mass in the interesting eV range without an intermediate gauge symmetry breaking scale one has to use a method that either involves a carefully chosen Yukawa coupling to an $\mathrm{SU}(2)_R$ triplet, which only arises for particular non-standard Kac-Moody levels, or non-renormalizable operators with $\mathrm{SU}(2)_R$ doublets [11]. Unlike in the case of an intermediate scale, both these methods require abandoning the attractice $b-\tau$ unification hypothesis except in the case of the $\mathrm{SU}(2)_R$ doublets and high $\tan\beta \sim m_t/m_b$ which requires greater tuning of the Higgs potential parameters and may also require M_G scale D-terms.

How the one loop MSSM beta functions have to be changed at an intermediate scale to increase the scale of gauge unification has been studied in a recent paper [12]. It is found that there exist only a few acceptable solutions with a single cleanly defined intermediate scale far below the unification scale. In fact if we demand an intermediate scale as mentioned in the above paragraph, the necessary relative changes in the beta functions of the MSSM are given as follows:

$$\Delta b_2 - \Delta b_1 = 2\,,\ \Delta b_3 - \Delta b_2 = 1\,, \tag{2}$$

where the hypercharge has been normalized in the standard GUT manner and $b_i = -2\pi\partial\alpha_i^{-1}/\partial\ln\mu$.

In this paper the additional field content we choose at the scale M_I is as follows: the additional vector representation fields necessary to complete the $\mathrm{SU}(2)_L\times\mathrm{SU}(2)_R\times\mathrm{SU}(4)_C$ symmetry, 2 copies of the chiral fields $H=(1,2,4)\equiv(\bar{u}^c_H,\bar{d}^c_H,\bar{E}^c_H,\bar{N}^c_H)$ and $\bar{H}=(1,2,\bar{4})\equiv$

$(u^c_H, d^c_H, E^c_H, N^c_H)$, and chiral singlets $S = (1,1,1)$ which are necessary for the right-handed neutrinos N^c_i to acquire large Majorana masses. We also add a chiral field $D = (1,1,6)$ to make all the non-MSSM Higgs modes massive along with two copies of chiral fields $\Phi = (2,2,1)$, which contain the MSSM Higgs. There are of course the usual 3 MSSM matter generations that include right-handed neutrinos $F = (2,1,4) \equiv (u,d,\nu,e)$ and $\bar{F} = (1,2,\bar{4}) \equiv (u^c, d^c, N^c, e^c)$. The $\mathrm{SU}(2)_R \times \mathrm{SU}(4)_C$ gauge symmetry is broken to the $\mathrm{U}(1)_Y \times \mathrm{SU}(3)_c$ by $\langle H \rangle = \left\langle \bar{N}^c_H \right\rangle$, $\left\langle \bar{H} \right\rangle = \langle N^c_H \rangle \sim M_I$. This causes 9 of the 21 gauge fields to become massive. The fields in H and $\bar{H}$ which combine with the corresponding components of the gauge fields are u^c_H, $\bar{u}^c_H$, E^c_H, $\bar{E}^c_H$, and a linear combination of N^c_H and $\bar{N}^c_H$ orthogonal to that of hypercharge to make super gauge-Higgs multiplets with a common mass of the order of M_I. In such a case, it is easy to see that any number of copies of H , $\bar{H}$ that might exist and get VEVs of order M_I would not leave any massless modes except for the ones corresponding to the hypercharge generator since all of the Higgsinos corresponding to these modes acquire mass through gaugino-Higgsino-<Higgs> terms. The linear combination of N^c_H and $\bar{N}^c_H$ corresponding to the hypercharge generator gets mass of the order of M_I through the terms $\lambda_{H\bar{H}S} H\bar{H}S$. Note that we do not add terms like $M_{H\bar{H}} H\bar{H}$ in the superpotential which cause the breaking of SUSY via F-terms. The presence of R symmetry [13] and the non-existence of a VEV for S could, for example, forbid these terms. The chiral fields d_H and $\bar{d}^c_H$ in these representations become massive with the help of the field $D = (1,1,6) \equiv (d^c_D, \bar{d}^c_D)$. This causes d_H, $\bar{d}^c_H$, d^c_D, and $\bar{d}^c_D$ to all get mass of order M_I through the superpotential terms $\lambda_{HHD} HHD$ and $\lambda_{\bar{H}\bar{H}D} \bar{H}\bar{H}D$ when H and $\bar{H}$ gets VEVS. Note that to avoid rapid proton decay, we also need to impose a symmetry on the superpotential (for example PQ symmetry as discussed later) that forbids terms of the type FFD and $\bar{F}\bar{F}D$ unless the couplings are extremely small.

The existence of the field D and S are crucial to make all the Higgs modes massive. As a matter of fact, in a previous paper Ref. [6] the field content without D and one of the bidoublets have been used to raise the unification scale. This field content also satisfies Eq.(2). (Note that in SO(10): $10 \rightarrow (1,1,6) + (2,2,1)$.) However the choice of this minimal

field content is problematic in practice as we have already pointed out that all the non-MSSM Higgs modes do not become massive at the breaking scale M_I, and hence the gauge coupling renormalization group equations (RGEs) are modified beneath the scale M_I. This model also suggests complete third generation Yukawa coupling unification at the intermediate scale, and hence requires large $\tan\beta$, due to the existance of only one bidoublet Higgs. Instead of the second bidoublet, we could have 2 copies of fields transforming as $(2,1,1)+(1,2,1)$. As a matter of fact, in string derivations of $\mathrm{SU}(2)_L\times\mathrm{SU}(2)_R\times\mathrm{SU}(4)_C$ models the exotic representations $(2,1,1)$, $(1,2,1)$, and $(1,1,4)+(1,1,\bar{4})$ tend to occur. Since the field content without D and S satisfy Eq. (2), the constraints on the additional field content is given by :

$$n_D \geq 1\,,\; n_D + n_4 = (n_\Phi - 1) + \frac{1}{2}n_2\,, \tag{3}$$

where n_D is the number of copies of fields transforming as $(1,1,6)$, n_Φ is the number of fields transforming as $(2,2,1)$, n_4 is the number of copies of $(1,1,4)+(1,1,\bar{4})$, and n_2 is the number of copies of $(2,1,1)+(1,2,1)$, which are all to be given mass of order M_I. Of course, Eq. (2) gives no constraint on n_S, the number of copies of singlet $(1,1,1)$ fields. It may be of interest to note that, for example, the first string derived version of the model found in Ref. [4] has $n_D=4$, $n_4=1$, $n_\Phi=4$, $n_2=10$, along with the necessary 2 copies of $H+\bar{H}$ ($N_H=2$) at the string scale and 3 generations of $F+\bar{F}$ ($N_F=3$), as well as several $\mathrm{SU}(2)_L\times\mathrm{SU}(2)_R\times\mathrm{SU}(4)_C$ singlets. Some of these $\mathrm{SU}(2)_L\times\mathrm{SU}(2)_R\times\mathrm{SU}(4)_C$ may acquire VEVs near the string scale which can break additional U(1) symmetries and may make some fields super heavy.

We now discuss how a low $\tan\beta$ scenario may be implimented in the model. To allow for the possibility of low $\tan\beta$, one needs the MSSM Higgs doublets ϕ_u and ϕ_d to not come primarily from the same bidoublet Φ_i. If we assume two bidoublets Φ_1 and Φ_2, no difficulty would exist in the breaking of the $\mathrm{SU}(2)_L\times\mathrm{SU}(2)_R\times\mathrm{SU}(4)_C$ model to the MSSM if one bidoublet was to remain light at M_I and the other was to be given mass of order M_I. However here we want one linear combination, from the two bidoublets, of down (or up) type Higgs superfield $\mathrm{SU}(2)_L$ doublets to remain massless and the other combination to

have mass of order M_I. (Otherwise, we would spoil the gauge coupling RGE analysis, which is the primary motivation for this model.) This can be accomplished through a modification of a method that has been used in conventional SO(10) GUTs. [14]. Consider adding to the model a pair of fields H_L and $\bar{H}_L$ transforming as $(2,1,4)$ and $(2,1,\bar{4})$, respectively, and also increasing the number of $H,\bar{H}$ pairs to be $N_H = 3$ so that Eqn.(2) is still satisfied. For simplicity, we discuss the general case and will not refer to specific choices of any PQ charges for the additional field content of this paragraph. If explicit mass terms for these fields are originally forbidden, they can be generated at the intermediate scale to be of order M_I through terms such as $\lambda_{H_L} H_L \bar{H}_L S_i$ where S_i gets a VEV at the scale M_I in the manner as previously discussed. In the limit of neglecting weak scale masses, the existence of the superpotential terms $W_D = \lambda_1 \bar{H} H_L \phi_1 + \lambda_2 H \bar{H}_L \phi_2$ and a symmetry (for example, the superpotential can be assumed to be invariant under the transformation: $H_L = -H_L$, $\phi_1 = -\phi_1$ and $S_i = -S_i$) forbidding the terms $\lambda_1 \bar{H}_L H \phi_1 + \lambda_2 H_L \bar{H} \phi_2$ would lead to the following form for the mass matrix for the $\mathrm{SU}(2)_L$ doublets:

$$M_D = \begin{pmatrix} \lambda_{H_L} v_{S_i} & \lambda_1 v_H & 0 \\ \lambda_2 v_{\bar{H}} & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \tag{4}$$

written in a basis where the rows stand for $(\bar{H}_{L_{\phi_u}}, \phi_{1u}, \phi_{2u})$ and the columns for $(H_{L_{\phi_d}}, \phi_{1d}, \phi_{2d})$ in obvious notation and all VEVs are of order M_I. This matrix naturally will have two large eigenvalues of order M_I, and has one massless eigenvalue which is composed of $\phi_{1d} = \phi_d$ and $\phi_{2u} = \phi_u$ and serves as the MSSM Higgs. Note that had D-parity not been broken in the model, we would not have had to have added any additional field content.

The gauge couplings in this model have the following one-loop beta functions:

$$b_i^{224} = \begin{pmatrix} -6 \\ -6 \\ -12 \end{pmatrix} + n_F \begin{pmatrix} 2 \\ 2 \\ 2 \end{pmatrix} + n_\Phi \begin{pmatrix} 1 \\ 1 \\ 1 \end{pmatrix} + n_D \begin{pmatrix} 1 \\ 1 \\ 1 \end{pmatrix} + N_{H_L} \begin{pmatrix} 0 \\ 4 \\ 2 \end{pmatrix} + N_H \begin{pmatrix} 4 \\ 0 \\ 2 \end{pmatrix}, \tag{5}$$

and the following two-loop beta functions:

$$b_{ij}^{224} = \begin{pmatrix} -24 & 0 & 0 \\ 0 & -24 & 0 \\ 0 & 0 & -96 \end{pmatrix} + N_F \begin{pmatrix} 14 & 0 & 15 \\ 0 & 14 & 15 \\ 3 & 3 & 31 \end{pmatrix} + n_\Phi \begin{pmatrix} 7 & 3 & 0 \\ 3 & 7 & 0 \\ 0 & 0 & 0 \end{pmatrix} + n_D \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 18 \end{pmatrix} \quad (6)$$
$$+ N_{H_L} \begin{pmatrix} 0 & 0 & 0 \\ 0 & 28 & 30 \\ 0 & 6 & 31 \end{pmatrix} + N_H \begin{pmatrix} 28 & 0 & 30 \\ 0 & 0 & 0 \\ 6 & 0 & 31 \end{pmatrix},$$

where $i = SU(2)_R, SU(2)_L, SU(4)_C$ respectively in the matrices, N_{H_L} is the number of field pairs(to be used for the low $\tan\beta$ scenario) transforming as $(2,1,4) + (2,1,\bar{4})$, $N_F = 3$ always, and we have left out the contributions from exotic representations since they can be easily calculated and we will not use them in our examples here. In Table 1, we show some sample gauge coupling unification results for the model in the case of $n_\Phi = 2$, $n_D = 1$ and $N_H = 2$ which requires complete third generation Yukawa coupling unification at the scale M_I and which we refer to as the high $\tan\beta$ scenario, for different values of $\alpha_s(M_Z)$ and the effective SUSY scale M_S with $\sin^2\theta_W = 0.2321$, $\alpha(M_Z) = 1/127.9$, $M_I = 10^{12}$ GeV, and the top quark mass $m_t \approx 180$ GeV. For the case that does not require $\lambda_t = \lambda_b$ at M_I, which we refer to as the low $\tan\beta$ scenerio, we use $n_\Phi = 2$, $n_D = 1$, $N_H = 3$ and $N_{H_L} = 1$. We display M_G, which we compare with the string scale prediction from Eq. (1). We do not include string threshold effects as we do not know the entire field content near the string scale, however we note from Ref. [2] that the effect of these thresholds in this model should in general tend towards having the bennificial effect of further reducing the small discrepancy between M_G and $M_{\rm string}$.

We note that the appearance of the intermediate breaking scale can occur through a single parameter in the singlet sector of the model. For example, consider an R symmetry invariant superpotential

$$W = \left(\sum_{i,j=1}^{2} \lambda_{ij} H_i \bar{H}_j - r \right) S_0 + ..., \quad (7)$$

where r is of order M_I and S_0 has no VEV. It would then be the most natural case for the VEVs acquired by all four fields to be of similar order. This mechanism can easily

be extended to link the breaking of a PQ-symmetry with the breaking of the intermediate gauge symmetry. For example, suppose a PQ-symmetry exists with $F, \bar{F}$ having a PQ charge of 1, bidoublet(s) Φ_i which contain the MSSM Higgs doublets and has (have) PQ charge -2, that the fields $H, \bar{H}$ have no or opposite PQ charges, singlets S_0, S_2, S_{-2} exist with the subscript denoting the PQ charge, and that R-symmetry prevents $M_{2,-2}S_2S_{-2}$ mass term from existing in the superpotential, then the following superpotential is possible:

$$W = \frac{\lambda_{\Phi_{ij}}}{M_{\rm Pl}}\Phi_i\Phi_j S_2 S_2 + \left(\lambda_S S_{-2}S_2 + \sum_{i,j=1}^{2}\lambda_{ij}H_i\bar{H}_j - r\right)S_0 + ... \tag{8}$$

where we have allowed a non-renormalizable term [15] so as to not need to fine-tune the μ parameter to a very large value or $\lambda_{\Phi_{ij}}$ to a tiny value. We observe that we can choose the PQ charge of D to be 0 for example so as to forbid terms like FFD and $\bar{F}\bar{F}D$ that would cause the rapid proton decay.

We note that in the model we are discussing there are no SU$(2)_R$ triplet fields, therefore one must rely on an extended version of the seesaw mechanism [16]. This mechanism in this scenario is described by the following 3×3 mass matrix for ν_a, N^c_a, and singlets S_a with $a = 1, 2, 3$:

$$M_\nu = \begin{pmatrix} 0 & \lambda_{(u)}v_u & 0 \\ \lambda^T_{(u)}v_u & 0 & fv_I \\ 0 & f^T v_I & M \end{pmatrix} \tag{9}$$

where $\lambda_{(u)}$, f, and M are 3×3 matrices, v_I is a VEV of order M_I, and v_u is the electroweak breaking VEV of the MSSM Higgs doublet H_u. Ignoring intergenerational mixing by pretending $\lambda_{(u)}$, f, and M are diagonal, then the mass of the ath light Majorana neutrino is given by $m_{\nu_a} \sim \left(\lambda_{(u)}v_u\right)^2 M_a/f_a^2 v_I^2$. To get back the usual seesaw relation and have m_{ν_τ} be in the eV range, we need $M_a \sim M_I$. This can be done by adding to the field content of the previous paragraph three generations of singlets transforming as S_{-1} and at least one S_1, assuming once again from R-symmetry that the explicit mass terms $M_{S_1,S_{-1}}S_1S_{-1}$ are forbidden in the SU$(2)_L\times$SU$(2)_R\times$SU$(4)_C$ superpotential, and giving $H_i, \bar{H}_j$ PQ charges of 0. All this allows the terms $\lambda_{ij}\bar{F}_iH_jS_{-1}$, $\lambda_{-1,-1,2}S_{-1}S_{-1}S_2$, and $\lambda_{-1,1,0}S_{-1}S_1S_0$ to be in the

superpotential, which at the intermediate scale would give S_{-1} both a mass and a VEV of order M_I and give the desired size to the entrees in the seesaw mixing matrix.

Lastly, we consider a signal of the model. Recently, GUT scale physics have been shown to be significant sources of lepton flavor violation in SUSY grand unification with a universal SUSY soft breaking boundary condition appearing near the reduced Planck scale [17–19]. However the parameter space where these signals could be observed is somewhat constrained by the experimental constarints of $b \to s\gamma$ [20]. More recently, it has also been shown that even if the universal boundary condition is taken at the GUT scale, that an intermediate gauge symmetry breaking can also be a significant source of lepton violation [22]. We now show that the rate of $\mu \to e\gamma$ can be within two orders of magnitude of experiment in the model discussed in this letter. For some parameter space, it can be above the experimental limit. In $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{SU}(4)_C$ gauge symmetry, the quarks and leptons are unified. Hence, the τ-neutrino Yukawa coupling is the same as the top Yukawa coupling. Through the RGE's, the effect of the large τ-neutrino Yukawa coupling is to make the third generation sleptons lighter than the first two generations, thus mitigating the GIM cancellation in one-loop leptonic flavor changing processes which involve virtual sleptons.

The superpotential terms which will be responsible for giving the SM fermion masses have the following form :

$$W_Y = \lambda_{\mathbf{F_u}} \mathbf{F} \mathbf{\Phi_2} \bar{\mathbf{F}} + \lambda_{\mathbf{F_d}} \mathbf{F} \mathbf{\Phi_1} \bar{\mathbf{F}}, , \tag{10}$$

where all group and generation indices have been suppressed. Φ_1 and Φ_2 are the two bidoublets. We have assumed that Φ_2 contains the MSSM Higgs doublet which gives masses to the up quarks and Dirac masses for the neutrinos and Φ_1 contains the doublet which gives masses to the down quarks and the charged leptons. Now we give the RGEs for the soft SUSY breaking parameters which we need for the intermediate gauge symmetry . First of all, there are gaugino masses M_i corresponding to each g_i. Secondly, corresponding to each tri-linear superpotential coupling $\lambda_{\mathbf{i}}$ there is a tri-linear scalar term with the coupling $\mathbf{A_i} \lambda_{\mathbf{i}}$ at M_G. Finally there are soft scalar mass terms m_i^2 for each of the the fields $F_{L,R}$, and $\Phi_{1,2}$.

$$\mathcal{D}\lambda^2_{F_{ug}} = -\sum_i c_i^{(\lambda_F)} g_i^2 + (4 + 4\delta_{g3})\,\lambda^2_{F_t}, \tag{11}$$

$$\mathcal{D}\lambda^2_{F_{dg}} = -\sum_i c_i^{(\lambda_F)} g_i^2 + 4\delta_{g3}\lambda^2_{F_t}, \tag{12}$$

$$\mathcal{D}M_i = b_i g_i^2 M_i, \tag{13}$$

$$\mathcal{D}A_{F_{ug}} = \sum_i c_i^{(\lambda_F)} g_i^2 M_i + (4 + 4\delta_{g3})\,\lambda^2_{F_t} A_{F_t}, \tag{14}$$

$$\mathcal{D}A_{F_{dg}} = \sum_i c_i^{(\lambda_F)} g_i^2 M_i + 4\delta_{g3}\lambda^2_{F_t} A_{F_t}, \tag{15}$$

$$\mathcal{D}m^2_{F,\bar{F}} = -\sum_i c_i^{\left(F,\bar{F}\right)} g_i^2 M_i^2 + 2\lambda^2_{F_t} X\delta_{g3}, \tag{16}$$

$$\mathcal{D}m^2_{\Phi_1} = -\sum_i c_i^{(\Phi)} g_i^2 M_i^2, \tag{17}$$

$$\mathcal{D}m^2_{\Phi_2} = -\sum_i c_i^{(\Phi)} g_i^2 M_i^2 + 4\lambda^2_{F_t} X, \tag{18}$$

where g refers to generation and i refers to the gauge,

$$c^{(\lambda_F)} = \left(3, 3, \frac{15}{2}\right),\ c^{(F)} = \left(3, 0, \frac{15}{2}\right), c^{(\bar{F})} = \left(0, 3, \frac{15}{2}\right), c^{(\Phi)} = (3, 3, 0)\,,$$
$$X \equiv m^2_F + m^2_{\bar{F}} + M^2_{\Phi_2} + A^2_{F_t},$$

and we have used

$$\mathcal{D} \equiv \frac{16\pi^2}{2}\frac{d}{dt},$$

where $t = \ln(\mu/\mathrm{GeV})$ with μ being the scale. At the scale M_G, we assume a universal form to the soft SUSY breaking parameters i.e. all gaugino masses $M_i(M_G) = m_{\frac{1}{2}}$, all tri-linear scalar couplings $A_i(M_G) = A_0$, and all soft scalar masses $m_i^2(M_G) = m_0^2$. At the scale M_I, we match the intermediate gauge symmetry breaking effective theory parameters with the MSSM parameters in the usual fashion. We run all the RGEs according to the MSSM [21] down to the top scale. Details of $\mu \to e\gamma$ with an intermediate symmetry are presented in Ref. [22]. As an example, we show the specific case of the universal gaugino masses $m_{1/2} = 145$ GeV and universal tri-linear soft breaking parameter $A_0 = 0$ at the unification scale, $\alpha_s(M_Z) = 0.119$, $M_G = 10^{17.87}$ GeV, $\alpha_G = 1/12.4$, $m_t = 176$ GeV and $m_b = 4.35$ GeV, and with minimal field content for low $\tan\beta = 2$ in Fig. 1. We have plotted the function

$$l_r \equiv \mathrm{Log}_{10}\left(\frac{B}{B_{\mathrm{exp}}}\right), \tag{19}$$

where B is the predicted $\mu \to e\gamma$ branching ratio and $B_{\rm exp} = 4.9 \cdot 10^{-11}$ being the experimental 90 % confidence limit upper bound on the branching ratio. With $A_0 = 0$, A_i is always negative at the weak scale. Consequently we find that $\mu < 0$ gives a greater branching ratio.

In summary, we have discussed a model that allows gauge coupling unification in the vicinity if the string scale and can have an intermediate $SU(2)_L \times SU(2)_R \times SU(4)_C$ gauge symmetry breaking scale M_I of order 10^{12} GeV, which is useful for producing a τ-neutrino with a mass of a few eV and solving the strong CP problem via a PQ symmetry whose breaking produces a harmless axion. We have shown that a range of field content is allowed by the model as given by Eq. (3) to satisfy Eqn. (2) which predicts $M_G/M_I \approx 10^{6\pm 2}$ GeV, and we have shown that Eqn. (1) which gives the string prediction between unification scale mass and gauge couplings are approximately satisfied for some choices of field content. We have also discussed what $SU(2)_L \times SU(2)_R \times SU(4)_C$ singlets and global U(1) charges may be useful to take advantage of $M_I \sim 10^{12}$ GeV and make a single parameter in the singlet sector responsible for the scale M_I.

Recently, an intermediate scale at $\sim 10^{12}$ GeV has been advocated [23] to produce monopoles that would explain the high energy cosmic ray spectrum. Since our model factors into U(1) group at M_I, it would be natural for such monopoles to arise in this model. It may be of interest to see if such monopoles satisfy the requirements of relic abundance required for the suggested mechanism.

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Table captions

Table 1 : The gauge unification scale M_G and string scale $M_{\rm string}$ (as predicted in Eqn. (1)) are shown for different values of $\alpha_s(M_z)$ and effective SUSY scale M_S in both the low and the high $\tan\beta$ models described in the text.

Figure captions

Fig. 1 : $l_r \equiv Log_{10}(B/B_{\rm exp})$ is plotted as a function of of the universal at M_G scale soft mass m_0.

The solid line corresponds to $\mu > 0$, while the dashed line corresponds to $\mu < 0$.

$\lambda_{F_{t_G}} = 1.25$ and $m_{1/2} = 140$ GeV for both lines.

$\tan\beta$	$\alpha_s(M_Z)$	$M_s(GeV)$	$1/\alpha_G$	$M_G(GeV)$	$M_{\rm string}(GeV)$
high	0.1258	175	20.34	$10^{18.26}$	$10^{17.74}$
high	0.1187	10^3	22.10	$10^{17.92}$	$10^{17.72}$
low	0.1192	175	10.65	$10^{17.83}$	$10^{17.84}$
low	0.1144	10^3	12.85	$10^{18.05}$	$10^{17.88}$

FIGURES

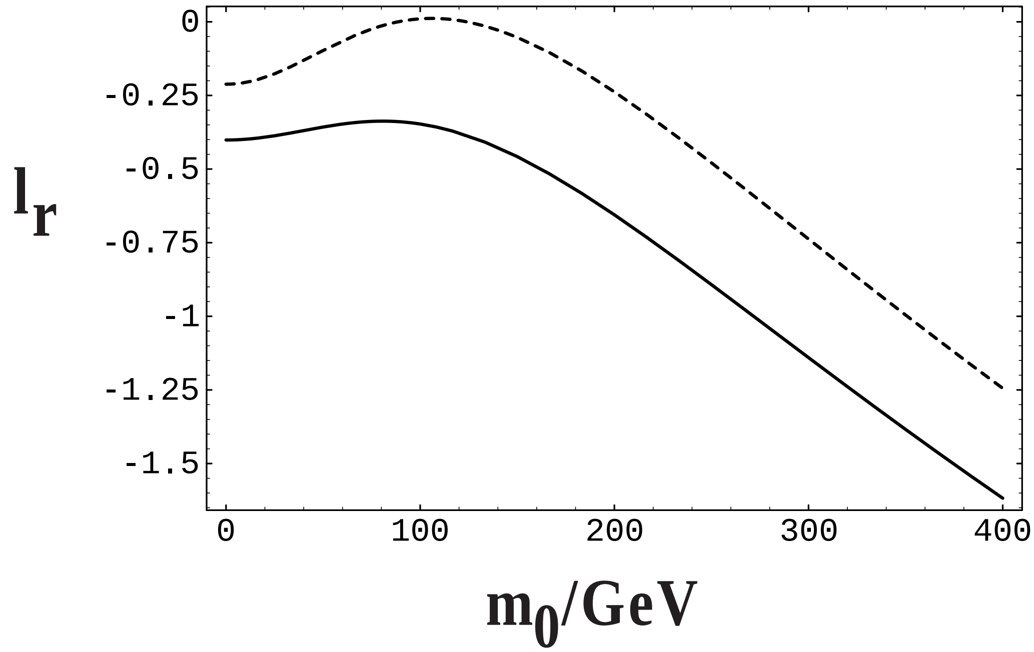

Fig. 1